Azimuthal asymmetry as a new handle on σ_L/σ_T in diffractive DIS

N.N. Nikolaev[a,b], A.V. Pronyaev[c], B.G. Zakharov[b]

[a] *IKP(Theorie), FZ Jülich, 5170 Jülich, Germany.*

[b] *D. Landau Institute for Theoretical Physics, GSP-1, 117940, ul. Kosygina 2, Moscow V-334, Russia.*

[c] *Virginia Polytechnic Institute and State University, Blacksburg, VA 24061-0435, USA.*

We propose a new method of the determination of $R^D = \sigma_L^D/\sigma_T^D$ from the dependence of the diffractive cross section on the azimuthal angle between the electron scattering and proton scattering planes. The method is based on our finding of the model independence of the ratio of the LT interference and transverse diffractive structure functions. The predicted azimuthal asymmetry is substantial and can be measured at HERA. We show that the accuracy of our reconstruction of R^D is adequate for a reliable test of an important pQCD prediction of $R^D \gtrsim 1$ for large β.

The ratio of (L) longitudinal and (T) transverse cross sections, $R = \sigma_L/\sigma_T$, is a much discussed test of mechanisms of deep inelastic scattering (DIS). The QCD theory of diffractive DIS (DDIS) $ep \to e'p'X$ predicts [1] an unprecedented dominance of higher twist σ_L^D over leading twist σ_T^D in a broad range of $\beta \gtrsim 0.9$ (Hereafter the superscript D stands for "diffractive" and the electron inelasticity y, the $\gamma^* p$ c.m.s. energy W, Q^2, $x = Q^2/(Q^2+W^2)$, the mass M of the diffractive system X, $\beta = Q^2/(Q^2+M^2)$ and $x_{\mathbf{IP}} = x/\beta$ are the standard DDIS variables). This must be contrasted to leading twist σ_L in small-x inclusive DIS where the theory predicts [2] $R = \sigma_L/\sigma_T \sim 0.2$-$0.3$ in agreement with the experiment [3]. The pQCD calculations of σ_L^D describe the experimental data for large β very well [4] and higher twist σ_L^D found in [1] has become a part of modern parameterizations of DDIS structure functions (SF's) [5]. However, a direct measurement of R^D requires a variable electron/proton energy runs which are not foreseen in the near future at HERA. Here below we show how this experimental limitation can be circumvented by the measurement of the LT interference SF F_{LT}^D.

The major point behind our proposal * is that Leading Proton Spectrometers of ZEUS and H1 opened an access to the transverse momentum $\vec{\Delta}$ of the recoil proton p' and an azimuthal angle ϕ between the electron scattering and proton production planes. The incident γ^* beam is a mixture of T and L photons and an interference of diffractive transitions $\gamma^*_{T,L} \to X$ leads to the $\cos\phi$ dependence of the observed cross section. Incidentally, such an LT interference is the s-channel helicity non-conserving (SCHNC) effect. There has been much discussion of different aspects of a related LT interference in exclusive electroproduction [7] and quasielastic scattering $A(e,e'p)$ [8]. The specific issue which we address in this paper is whether it is really possible to deduce $R^D = \sigma_L^D/\sigma_T^D$ from the experimentally measured $\sigma_2^D = \sigma_L^D + \sigma_T^D \propto \sum_X \left\{|A_L(\gamma^* \to X)|^2 + |A_T(\gamma^* \to X)|^2\right\}$ and $\sigma_{LT}^D \propto \mathrm{Re}\sum_X A_L^*(\gamma^* \to X) A_T(\gamma^* \to X)$ and whether the so reconstructed R^D will be model independent and under the control of perturbative QCD?

The principal point is as follows. For unpolarized electrons the differential cross-section of DDIS can be decomposed as

$$Q^2 y \frac{d\sigma(ep \to ep'X)}{dQ^2 dy dM^2 d\Delta^2 d\phi} = \frac{\alpha_{em}}{2\pi^2}\left[\left(1 - y + \frac{y^2}{2}\right)\frac{d\sigma_T^D}{dM^2 d\Delta^2} + (1-y)\frac{d\sigma_L^D}{dM^2 d\Delta^2} + (1-y)\frac{d\sigma_{TT}^D}{dM^2 d\Delta^2}\cdot\cos 2\phi + \left(1-\frac{y}{2}\right)\sqrt{1-y}\,\frac{d\sigma_{LT}^D}{dM^2 d\Delta^2}\cdot\cos\phi\right] =$$
$$\frac{\alpha_{em}}{2\pi^2}\cdot\left(1 - y + \frac{y^2}{2}\right)\cdot\frac{d\sigma_2^D}{dM^2 d\Delta^2}\left\{1 - \frac{y^2 R^D}{2(1-y)+y^2} + \frac{2(1-y)A_{TT}}{2(1-y)+y^2}\cos 2\phi + \frac{(2-y)\sqrt{1-y}A_{LT}}{2(1-y)+y^2}\cos\phi\right\}. \quad (1)$$

The experimental measurement of the asymmetry $\propto \cos\phi$ amounts to the determination of

$$A_{LT} = \frac{F_{LT}^D}{F_T^D + F_L^D} = \frac{\rho_{LT}}{1+R^D}. \quad (2)$$

(for the sake of simplicity we focus on $y \ll 1$ which is the standard experimental cutoff). If $\rho_{LT} = F_{LT}^D/F_T^D$ were known, then inversion of (2) would yield

$$R^{D(4)} = \frac{\rho_{LT}}{A_{LT}} - 1. \quad (3)$$

Below we argue for $\beta \gtrsim 0.9$ of the interest the theoretical evaluations of ρ_{LT} are model-independent and the inversion (3) offers a viable test of the pQCD prediction $R^D \gtrsim 1$ from the experimentally measured azimuthal asymmetry A_{LT}.

The microscopic QCD mechanism of DDIS at $\beta > 0.9$ is an excitation of $q\bar{q}$ Fock states of the photon (we focus on continuum $M^2 \gg 4m_f^2$). The kinematical variables are shown in fig. 1, quark and antiquark carry a fraction z and $1-z$ of the photon's momentum and

$$M^2 = Q^2\frac{1-\beta}{\beta} = \frac{\vec{k}^2 + m_f^2}{z(1-z)}. \quad (4)$$

*The preliminary results from this study have been reported elsewhere [6]

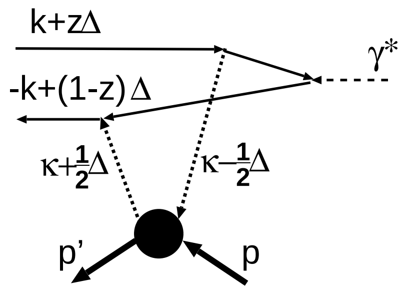

FIG. 1. The sample Feynman diagram for diffraction excitation of the quark-antiquark Fock state of the virtual photon.

In our calculation of SCHNC σ^D_{LT} which is a new result of this work we rely heavily upon the theory of the s-channel helicity conserving (SCHC) DDIS developed in [1,4,9,10]. Extension [11] of this formalism to the $\vec{\Delta}\neq 0$ gives ($i = L, T, LT, TT, n_L = n_T = 0, n_{LT} = 1, n_{TT} = 2$)

$$\frac{d\sigma_i^D}{dM^2 d\Delta^2}\cdot\cos(n_i\phi) = \frac{\alpha_{em}}{12\pi}\sum_f e_f^2\int d^2\vec{k}\frac{k^2+m_f^2}{JM^4}\alpha_S^2(\overline{Q}^2)[H_i(z_+)+H_i(z_-)]\,, \quad (5)$$

where e_f is the quark charge in units of the electron charge, m_f is the quark mass, $z_\pm = \frac{1}{2}(1\pm J)$, $J^2 = 1-\frac{4(k^2+m_f^2)}{M^2}$ and the hard scale $\overline{Q}^2$ is defined below. For the SCHC cross sections we have the familiar results [9] $H_T = [1-2z(1-z)]\vec{\Phi}_1^2 + m_f^2\Phi_2^2$ and $H_L = 4z^2(1-z)^2Q^2\Phi_2^2$, whereas for the SCHNC cross sections we find

$$H_{LT} = 4z(1-z)(1-2z)Q(\vec{\Phi}_1\vec{t})\Phi_2\,, \quad (6)$$

$$H_{TT'} = 2z(1-z)[\vec{\Phi}_1^2 - 2(\vec{\Phi}_1\vec{t})^2]\,, \quad (7)$$

where $\vec{t}$ is a unit vector tangential to the (e, e') scattering plane and orthogonal to the $\gamma^* p$ collision axis. Diffractive amplitudes $\vec{\Phi}_1$ and Φ_2 are calculable [4,9,11] in terms of the lightcone wave function of the $q\bar{q}$ Fock state of the photon $\psi_2(z,\vec{k}) = 1/[\vec{k}^2+m_f^2+z(1-z)Q^2]$ (we also use $\vec{\psi}_1(z,\vec{k}) = \vec{k}\psi_2(z,\vec{k})$) and gluon density matrix of the proton $\mathcal{F}\left(x_{\bf IP},\vec{\kappa},\vec{\Delta}\right)$:

$$\Phi_j = \frac{1}{2\pi}\int\frac{d^2\vec{\kappa}}{\kappa^4}\mathcal{F}\left(x_{\bf IP},\vec{\kappa},\vec{\Delta}\right)[\psi_j(z,\vec{r}+\vec{\kappa}) + \psi_j(z,\vec{r}-\vec{\kappa}) - \psi_j(z,\vec{r}+\frac{\vec{\Delta}}{2}) - \psi_j(z,\vec{r}-\frac{\vec{\Delta}}{2})]\,, \quad (8)$$

and $\vec{r} = \vec{k} - \frac{1}{2}(1-2z)\vec{\Delta}$. The dependence of $\mathcal{F}(x_{\bf IP},\vec{\kappa},\vec{\Delta})$ on the variable $\vec{\Delta}\vec{\kappa}$ corresponds to the subleading BFKL singularities [12] and can be neglected for small $x_{\bf IP}$ of the interest. For small $\vec{\Delta}^2$ within the diffraction cone

$$\mathcal{F}(x,\vec{\kappa},\vec{\Delta}) = \frac{\partial G(x,\kappa^2)}{\partial\log\kappa^2}\exp(-\frac{1}{2}B_{3\bf IP}\vec{\Delta}^2)\,, \quad (9)$$

where $\partial G/\partial\log\kappa^2$ is the conventional unintegrated gluon structure function and the diffraction cone $B_{3\bf IP}\sim 6$ GeV^{-2} [13]. As usual [1,4,10], the dominant contribution to $\vec{\Phi}_1, \Phi_2$ comes from the leading log$\overline{Q}^2$ region of

$$\vec{\kappa}^2 \lesssim \overline{Q}^2 = \frac{\vec{k}^2+m_f^2}{1-\beta} \quad (10)$$

with the results

$$\vec{\Phi}_1 = \frac{2\vec{r}}{\overline{Q}^4}\left[\beta+\frac{m_f^2}{\overline{Q}^2}\right]G(x_{\bf IP},\overline{Q}^2)\,, \quad (11)$$

$$\Phi_2 = \frac{1}{\overline{Q}^4}\left[2\beta-1+\frac{2m_f^2}{\overline{Q}^2}\right]G(x_{\bf IP},\overline{Q}^2)\,. \quad (12)$$

For small $\vec{\Delta}$ the nonvanishing contribution to σ_{LT} comes from the term $\propto (1-2z)\vec{\Delta}$ in $\vec{r}$, so that $H_{LT}\propto(\vec{t}\vec{\Delta}) = \Delta\cdot cos\phi$ as expected from the SCHNC interference of the helicity-flip and helicity non-flip diffractive amplitudes. The typical contributions to the k^2 integrated cross section are proportional to integrals ($i = T, L, LT$)

$$\mathcal{J}_i^f(x_{\bf IP},\langle\overline{Q}^2\rangle_i) = (n+1)\int_0^{\frac{M^2}{4}-m_f^2}\frac{dk^2}{k^2+m_f^2}\times\left[\frac{m_f^2}{k^2+m_f^2}\right]^n\left[\alpha_S(\overline{Q}^2)G(x_{\bf IP},\overline{Q}^2)\right]^2 \quad (13)$$

where to the leading log $\langle\overline{Q}^2\rangle_i$ approximation and for heavy flavours $\mathcal{J}_i^f = \left[\alpha_s(\langle\overline{Q}^2\rangle_i)G(x_{\bf IP},\langle\overline{Q}^2\rangle_i)\right]^2$, for the case of light flavours see a discussion in [4]. We recall that for σ_T^D which is twist-2 we have $n\geq 1$ and the dominant contribution to σ_T^D comes from the aligned jet configurations, $k^2\sim m_f^2$. Consequently [4,10] the relevant average QCD hard scale $\langle\overline{Q}^2\rangle_T\sim\frac{m_f^2}{1-\beta}$ which is always large for heavy flavours and/or $1-\beta\ll 1$.

Our new finding is that σ_{LT} also is dominated by the aligned jet configurations, which can readily be checked from eqs. (5),(6),(11) and (12). Consequently, we expect

$$\langle\overline{Q}^2\rangle_{LT}\approx\langle\overline{Q}^2\rangle_T\,, \quad (14)$$

which is a basis for our conclusion on the model-independence of ρ_{LT}. Indeed, our result for the SCHNC DDIS SF $F_{LT}^{D(4)}$ reads

$$F_{LT}^{D(4)}(\Delta^2,x_{\bf IP},\beta,Q^2) = \frac{\Delta}{Q}\cdot\frac{2\pi}{9\sigma_{tot}^{pp}}\cdot 24\beta^4(1-\beta)(2-3\beta)\times\sum_f\frac{e_f^2}{m_f^2}\mathcal{J}_{LT}^f(x_{\bf IP},\langle\overline{Q}^2\rangle_{LT})\exp\left(-B_{LT}^f\Delta^2\right)\,. \quad (15)$$

(Incidentally, F_{LT}^D is twist-3). It must be compared to $F_T^{D(4)}$ of [4]:

$$F_T^{D(4)}(\Delta^2, x_{\mathbf{IP}}, \beta, Q^2) = \frac{2\pi}{9\sigma_{tot}^{pp}} \cdot \beta(1-\beta)^2(3+4\beta+8\beta^2) \times$$
$$\sum_f \frac{e_f^2}{m_f^2} \mathcal{J}_T^f(x_{\mathbf{IP}}, \langle \overline{Q}^2 \rangle_T) \exp\left(-B_T^f \Delta^2\right) . \quad (16)$$

Then, our finding of the proximity of the two hard scales (14) entails $\mathcal{J}_{LT}(x_{\mathbf{IP}}, \langle\overline{Q}^2\rangle_{LT}) \approx \mathcal{J}_T(x_{\mathbf{IP}}, \langle\overline{Q}^2\rangle_T)$ and for the dominant contribution from light flavour excitation we obtain the fully analytic approximation

$$\rho_{LT}(\beta,\Delta) = \chi_{LT}(\beta) \cdot \frac{\Delta}{Q} \cdot \exp(-[B_{LT} - B_T]\vec{\Delta}^2) =$$
$$\frac{24\beta^3\,(2-3\beta)}{(1-\beta)\,(3+4\beta+8\beta^2)} \cdot \frac{\Delta}{Q} \cdot \exp(-[B_{LT} - B_T]\vec{\Delta}^2)\,, \quad (17)$$

in which the *r.h.s.* depends on neither $x_{\mathbf{IP}}$ nor Q^2 (apart from the trivial kinematical factor Δ/Q). For $\beta \gtrsim 0.9$ at practically attainable $Q^2 \lesssim 100$ GeV^2 the open charm contribution to F_T^D and F_{LT}^D is negligible small, see [4] and here below.

Obviously the accuracy of our evaluation of R^D from eq. (3) hinges on the knowledge of diffraction slopes B_i ($i = T, L, LT$) and an accuracy of the analytic approximation (17). Given the parameterization of the gluon structure function of the proton, one can calculate the ratio ρ_{LT} from the first principles using eqs. (5)-(8) and compare such a numerical result with the analytic approximation (17).

Before venturing into this numerical experiment, we notice that for light flavours and moderately small $1-\beta \sim 0.05$-0.1 the both F_{LT} and F_T receive certain contribution from the soft-to-hard transition region of $\overline{Q}^2$. For $Q^2 \gtrsim Q_c^2$, where $Q_c^2(GRV) = 0.4, Q_c^2(CTEQ) = 0.49, Q_c^2(MRRS) = 1.25$ GeV^2, we can evaluate $\mathcal{F}(x_{\mathbf{IP}}, Q^2)$ using the standard GRV NLO [14], CTEQ 4LQ [15] and/or MRRS ($m_c = 1.5 GeV$) [16] parameterizations for the perturbative gluon SF of the proton and at soft Q^2 we apply to $\mathcal{F}(x_{\mathbf{IP}}, Q^2)$ soft-to-hard interpolation described in [4]. The minor improvement over the related interpolation of $G(x, Q^2)$ in [4] is that we constrain the extrapolation of perturbative $\mathcal{F}(x_{\mathbf{IP}}, Q^2)$ to have the Q^2 dependence as in eq. (14) of [4] with the infrared cutoff parameter $\mu_G = 0.75$ GeV [17]. This way we correctly impose the gauge-invariance driven cancellation of radiation of soft gluons by colorless protons. Then we apply to $\mathcal{F}(x_{\mathbf{IP}}, Q^2)$ the interpolation eq. (15) of [4] and request the same numerical result for the transverse diffractive SF F_T^D at the typical $x_{\mathbf{IP}} = 10^{-3}$ and $Q^2 = 100$ GeV^2 and $\beta = 0.7$ as found in [4], which fixes the nonperturbative soft contribution parameter: $C_B(GRV) = 1.7, C_B(CTEQ) = 1.86, C_B(MRRS) = 2.0$. For all three models the unintegrated gluon SF $\mathcal{F}(x_{\mathbf{IP}}, Q^2)$ builds up from zero on a typical scale $\overline{Q}_G^2 \sim 0.5$ GeV^{-2} and for light flavours it is $\overline{Q}_G^{-2}$ which supplants m_f^{-2} as a scale in normalization of diffractive SF's (15),(16). We checked that when $\beta \to 1$ starting from this reference $\beta = 0.7$ and the QCD scale Q_T^2 rises from the soft value to hard $Q_T^2 \to \frac{1}{4}Q^2$, the values of F_T^D for the three gluon densities diverge by $\lesssim$ 5-10% in conformity to the known slight divergence of the GRV, CTEQ and MRRS gluon densities. We checked also that the results for F_L^D for the three models agree within the same accuracy. In order to not have very heavy figures, we usually show the numerical results for the CTEQ gluons. All the numerical results will be shown for $x_{\mathbf{IP}} = 10^{-3}$ and $Q^2 = 100$ GeV^2 which are typical of $\beta \sim 1$ kinematics at HERA. The continuum calculations can be trusted for $M^2 \gtrsim 3$ GeV^2, beyond the exclusive resonance excitation, and we only consider $\beta < 0.97$.

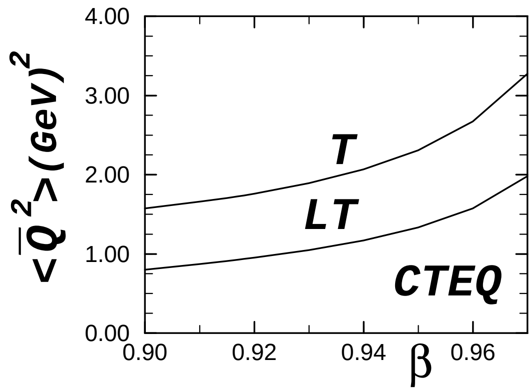

FIG. 2. The average hard scale $\langle\overline{Q}^2\rangle$ for light flavour contribution to the (T) transverse and LT interference diffractive structure functions evaluated for the CTEQ gluon SF of the proton.

The average value of the running hard scale (10) which enters the calculations of the σ_{LT} and σ_T can be estimated calculating the expectation values of $(\overline{Q}^2)^\gamma$. For heavy flavours and weak scaling violations in the gluon structure SF these moments can be evaluated analytically with the result $\langle\overline{Q}^2\rangle_T = \left(\frac{6}{5}\right)^2 \langle\overline{Q}^2\rangle_{LT}$ for $\gamma = \frac{1}{2}$, the inequality being due to a somewhat different contribution of terms (13) with $n=1$ and $n=2$ in σ_{LT} and σ_T. Numerical calculation for light flavours yields a comfortably large $\langle\overline{Q}^2\rangle = \langle(\overline{Q}^2)^\gamma\rangle^{\frac{1}{\gamma}}$ shown in fig. 2 for the typical $\gamma = \frac{1}{2}$, which confirms the expectation (14). The slight inequality $\langle\overline{Q}^2\rangle_T \approx 1.6\langle\overline{Q}^2\rangle_{LT}$ is similar to that for heavy flavours and does not affect our major argument.

The calculation of the diffraction slope $B_{L,T}$ for the SCHC structure functions $F_{L,T}^D$ is found in [11], here we cite our new result for B_{LT}:

$$B_{LT} = B_{3\mathbf{IP}} +$$
$$\frac{1}{20m_f^2}\frac{(1-\beta)(2+7\beta+12\beta^2-483\beta^3+672\beta^4)}{12\beta^2(2-3\beta)} \quad (18)$$

Here the second term is the contribution from the γ^*X

transition vertex and is a rigorous pQCD result because for heavy flavours $\langle\overline{Q}^2\rangle_{LT}$ is large for all β (barring the vicinity of $\beta=\frac{2}{3}$ in which $F_{LT}^{D(4)}(\Delta=0)$ vanishes and B_{LT} is ill defined). In the case of light flavours the scale of this component of B_{LT} (and B_T too [11]) is rather set by $\overline{Q}_G^{-2}$ than m_f^{-2}. The variations of diffraction slopes from CTEQ to GRV to MRS parameterizations do not exceed ~ 5 per cent and are not shown in fig. 3.

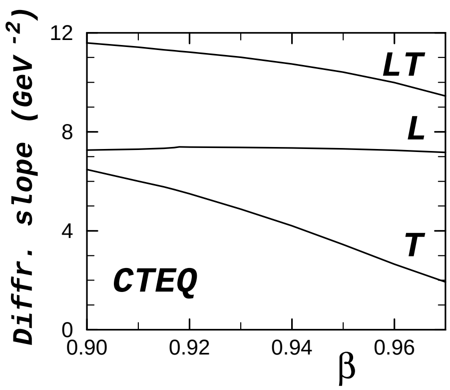

FIG. 3. Our predictions for the diffraction slopes B_T, B_L, B_{LT} of the transverse, longitudinal and LT interference diffractive SF's for the CTEQ gluon SF of the proton ($x_{\mathbf{I\!P}}=10^{-3}$, $Q^2=100$ GeV^2).

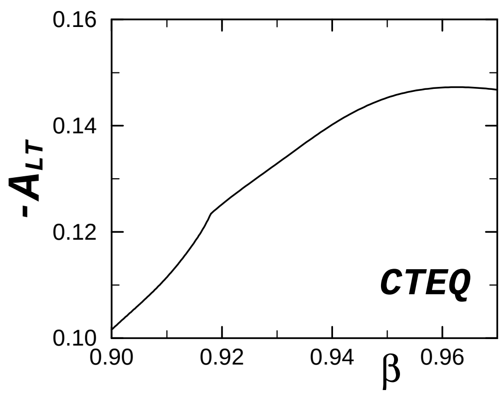

FIG. 4. Our predictions for the β dependence of the azimuthal asymmetry A_{LT} for the Δ-integrated cross sections evaluated using the diffraction slopes of fig. 3 ($x_{\mathbf{I\!P}}=10^{-3}$, $Q^2=100$ GeV^2, CTEQ gluon SF of the proton).

Evidently, a linear growth of azimuthal asymmetry with Δ stops and the asymmetry reaches its maximal value at (B_L is more relevant in the region of the dominance of F_L)

$$\Delta^2=\frac{1}{2(B_{LT}-B_L)}\sim(0.1-0.2)\ \mathrm{GeV}^2 \tag{19}$$

which lies comfortably within the acceptance of the LPS's at HERA. With sufficiently high statistics, one can study experimentally the full Δ dependence of the azimuthal asymmetry. In order to have the first impressions of the expected signal, hereafter we shall consider the Δ integrated quantities, in which case

$$\rho_{LT}=\frac{1}{2}\sqrt{\frac{\pi}{B_{LT}Q^2}}\cdot\frac{B_T}{B_{LT}}\cdot\chi_{LT}(\beta)\,, \tag{20}$$

$$R^D=R^D(\Delta=0)\frac{B_T}{B_L}\,. \tag{21}$$

The predicted azimuthal asymmetry A_{LT} is shown in fig. 4. It is substantial and within the reach of the LPS experiments at HERA. For $Q^2\sim 50$ GeV^2 for which $\beta\sim 0.95$ does still belong to the continuum, the expected asymmetry is $\approx 50\%$ larger. The cusp-like behaviour at $\beta=\beta_c=Q^2/(Q^2+4m_c^2)\approx 0.92$ is due to the open charm excitation at $\beta<\beta_c$.

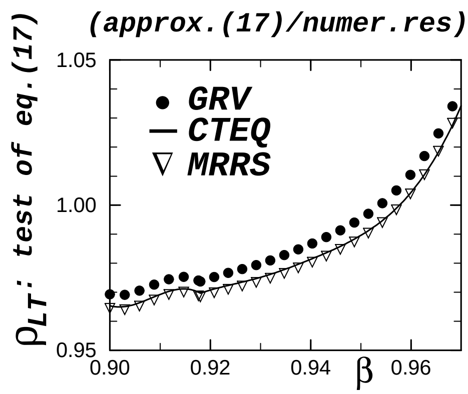

FIG. 5. Numerical test of the accuracy of the analytic formula (17). Shown is the ratio of approximation (17) for ρ_{LT} to the direct numerical evaluation of ρ_{LT} ($x_{\mathbf{I\!P}}=10^{-3}$, $Q^2=100$ GeV^2).

Now we are in the position to test the accuracy of the analytic approximation (17). In fig. 5 we show the ratio of ρ_{LT} given by (17) to the direct numerical evaluation of ρ_{LT}. A departure of their ratio from unity does not exceed 5 per cent which confirms the expected high accuracy of our analytic result (17). The cusp is a weak effect of the open charm threshold which has been neglected in (17) but included in the numerical evaluation of ρ_{LT}. The found variation from the GRV to CTEQ to MRRS gluons is negligible for the practical purposes.

The accuracy of reconstruction (3) of $R^D=\sigma_L^D/\sigma_T^D$ can be judged from fig. 6a. Here we show the ratio of R^D reconstructed from the numerically calculated azimuthal asymmetry A_{LT} on the basis of analytic approximation (17) for ρ_{LT} to the direct numerical result for R^D presented in fig. 6b (for the convenience of plotting, in fig. 6b we show the inverse quantity $1/R^D(\Delta=0)$). The reconstruction errors do dot exceed dozen per cent and this accuracy of our reconstruction of R^D is adequate for a reliable check of the pQCD predictions of large R^D at

$\beta \gtrsim 0.9$. We emphasize the Q^2 independence of $\langle\overline{Q}^2\rangle_{LT}$ and $\langle\overline{Q}^2\rangle_T$ by which ρ_{LT} does not depend on Q^2. Consequently, our method allows a reliable experimental determination of the Q^2 dependence of R^D and an unambiguous test of the higher twist nature of σ_L^D.

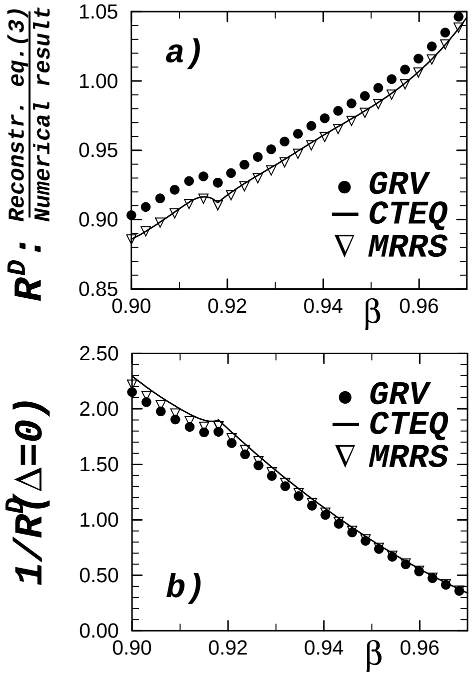

FIG. 6. a) A comparison of the $R^D = \sigma_L/\sigma_T$ reconstructed from azimuthal asymmetry of fig. 4 using eqs. (3) and approximation (17) for ρ_{LT} with the direct numerical evaluation of R^D for the GRV, CTEQ and MRRS gluon SF of the proton. b) Our numerical predictions for $R^D(\Delta = 0)$ in forward DDIS for the same set of gluon SF's of the proton ($x_{\mathbf{IP}} = 10^{-3}$, $Q^2 = 100$ GeV^2).

To summarize, we derived the s-channel helicity non-conserving LT interference diffractive structure function F_{LT}^D. It leads to a substantial dependence of the diffractive cross section on the azimuthal angle between the electron scattering and proton scattering planes which can be observed at HERA. We demonstrated a weak model dependence of the ratio F_{LT}^D/F_T^D for large β, which makes viable the reconstruction of the otherwise inaccessible $R^D = \sigma_L^D/\sigma_T^D$ from the experimentally measured azimuthal asymmetry. We checked that the accuracy of our method is adequate for the reliable test of the striking pQCD prediction of $R^D \gtrsim 1$ for large β.

The work of B.G.Z. has been supported partly by the INTAS grant 96-0597 and the work of A.V.P. was supported partly by the US DOE grant DE-FG02-96ER40994.

[1] Genovese M., N.Nikolaev and B.Zakharov, *Phys.Lett.* **B380**, 213 (1996).

[2] N. N. Nikolaev and B.G.Zakharov, *Phys. Lett.* **B327**, 149 (1994); **B327**, 157 (1994).

[3] H1 Collab., C. Adloff, S. Aid, M. Anderson, V. Andreev, B. Andrieu et al., *Phys. Lett.* **B393**, 452 (1997); E140X Collab., L.H. Tao, L. Andivahis, P. Anthony, R. Arnold et al. *Z. Phys.* **C70** 387 (1996); E143 Collab., K. Abe, T. Akagi, P.L. Anthony, R. Antonov et al., **SLAC-PUB-7927**, Aug 1998, hep-ex/9808028.

[4] Bertini M., M. Genovese, N.N. Nikolaev, A.V. Pronyaev and B.G. Zakharov, *Phys. Lett.* **B422**, 238 (1998).

[5] J. Bartels, J. Ellis, H. Kowalski and M. Wusthoff, **CERN-TH-98-67**, hep-ph/9803497; J. Bartels and C. Royon, hep-ph/9809344.

[6] N.N.Nikolaev and B.G.Zakharov, Phenomenology of Diffractive DIS, in: Deep Inelastic Scattering and QCD (DIS97), Proc. of 5th International Workshop, Chicago, IL USA, April 14-18, 1997, American Institute of Physics Proceedings No.407, edited by J.Repond and D.Krakauer, pp. 445-455; A.V.Pronyaev, The Forward Cone and L/T Separation in Diffractive DIS, hep-ph/9808432, To be published in the proceedings of 6th International Workshop on Deep Inelastic Scattering and QCD (DIS 98), Brussels, Belgium, April 4-8, 1998.

[7] J. Speth and V.R. Zoller, *Phys. Lett.* **B351**, 533 (1995) C.J. Bebek, A. Browman, C.N. Brown, K.M. Hanson et al. *Phys. Rev. Lett.* **37**, 1326 (1976).

[8] S. Boffi, C. Giusti and F.D. Pacati, *Phys. Rept.* **226**, 1 (1993); M. Holtrop, D. Jordan, T. McIlvain, R. Alarcon et al., nucl-ex/9709003.

[9] Nikolaev N.N. and B.G.Zakharov, *Phys. Lett.* **B332**, 177 (1994); *Z. Phys.* **C53**, 331 (1992).

[10] Genovese M., N.Nikolaev and B.Zakharov, *Phys. Lett.* **B378**, 347 (1996).

[11] N.N.Nikolaev, A.V.Pronyaev and B.G.Zakharov, JETP Lett. 68, 604 (1998).

[12] L.N.Lipatov, *Sov. Phys. JETP* **63** (1986) 904; L.N.Lipatov, in:*Perturbative Quantum Chromodynamics*, ed. by A.H.Mueller, World Scientific (1989); E.A.Kuraev, L.N.Lipatov and S.V.Fadin, *Sov. Phys. JETP* **44** (1976) 443; *Sov. Phys. JETP* **45** (1977) 199;

[13] J.Nemchik, N.N.Nikolaev, E.Predazzi, B.G.Zakharov and V.R.Zoller, *J. Exp. Theor. Phys.* **86**, 1054 (1998).

[14] M. Gluck, E. Reya and A. Vogt, *Z. Phys.* **C67**, 433 (1995).

[15] H.L. Lai, J. Huston, S. Kuhlmann, F. Olness et al., *Phys. Rev.* **D55**, 1280 (1997).

[16] A.D.Martin, R.G.Roberts, M.Ryskin and W.J.Stirling, *Eur. Phys. J.* **C2**, 287 (1998).

[17] N.N.Nikolaev and B.G.Zakharov, *Phys. Lett.* **B327** (1994) 149; *Phys. Lett.* **B327** (1994) 157.